Parity: from strong CP problem to dark matter, neutrino masses and baryon asymmetry

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We show that in an $SU(3)_c\times[SU(2)_L\times U(1)_Y]\times[SU(2)'_R\times U(1)'_Y]$ framework, the parity symmetry motivated by solving the strong CP problem without resorting to an axion can predict a dark matter particle with a mass around 302 GeV. This dark matter candidate can be directly detected in the presence of a $U(1)_Y\times U(1)'_Y$ kinetic mixing. Furthermore, our model can accommodate a natural way to simultaneously realize an inverse-linear seesaw for neutrino masses and a resonant leptogenesis for baryon asymmetry.

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I. INTRODUCTION

The most popular solution to the so-called strong CP problem is to introduce a continuous Peccei-Quinn [1] symmetry which predicts a pseudo Goldstone boson, the axion [1–4], and hence a dynamical strong CP phase. However, the axion has not been experimentally seen so far. Alternatively, we can consider certain discrete symmetries to suppress or remove the strong CP phase. For example, Babu and Mohapatra [5] proposed an $SU(3)_c\times SU(2)_L\times SU(2)_R\times U(1)_{B-L}$ left-right symmetric model [6] with parity symmetry to generate a tiny strong CP phase at two-loop level. Barr, Chang and Senjanović [7] then pointed out the Babu-Mohapatra scheme could be generalized in a relaxed $SU(3)_c\times[SU(2)_L\times U(1)_Y]\times[SU(2)'_R\times U(1)'_Y]$ framework. Such dark parity extension of the ordinary electroweak theory could result in some stable dark particles, among which the stable dark quark is not allowed to have a significant relic density. To dilute the stable dark quark in the early universe, one may expect an inflation with a reheating temperature much below the mass of the stable dark quark [7].

In this paper we shall demonstrate a novel $SU(3)_c\times[SU(2)_L\times U(1)_Y]\times[SU(2)'_R\times U(1)'_Y]$ model with some interesting implications of the parity symmetry motivated by solving the strong CP problem on the property of dark matter, the origin of neutrino masses and the generation of baryon asymmetry. In our model, after the dark electromagnetic symmetry is spontaneously broken, a dark up quark can have a fast decay into a dark electron and four ordinary fermions through the mediations of three colored scalars. Our model also contains three gauge-singlet fermions with small Majorana masses and an $[SU(2)_L\times SU(2)'_R]$-bidoublet Higgs scalar with seesaw-suppressed vacuum expectation value (VEV). The singlet fermions and the dark neutrinos can form three pairs of quasi-degenerate Majorana fermions to realize a resonant [8, 9] leptogenesis [10, 11] for baryon asymmetry. Meanwhile, the neutrino masses can have a form of inverse [12] and linear [13] seesaw [14, 15]. The dark electron can annihilate into the dark photon and then obtain a relic density to explain the dark matter puzzle. The dark matter relic density will only depend on the dark electron mass since the parity symmetry identifies the dark gauge couplings to the ordinary ones. We hence can determine the dark matter mass to be about 302 GeV from the measured dark matter relic density. In the presence of a $U(1)$ kinetic mixing, our dark matter particle can be verified by the ongoing and future dark matter direct detection experiments.

II. THE MODEL

The ordinary and dark scalars are denoted by

$$\begin{aligned}
\phi(\mathbf{1})(\mathbf{2},-\tfrac{1}{2})(\mathbf{1},0) &\leftrightarrow \phi'(\mathbf{1})(\mathbf{1},0)(\mathbf{2},-\tfrac{1}{2})\,,\\
\eta(\mathbf{1})(\mathbf{1},-\tfrac{1}{3})(\mathbf{1},0) &\leftrightarrow \eta'(\mathbf{1})(\mathbf{1},0)(\mathbf{1},-\tfrac{1}{3})\,,\\
\delta(\mathbf{3})(\mathbf{1},-\tfrac{1}{3})(\mathbf{1},0) &\leftrightarrow \delta'(\mathbf{3})(\mathbf{1},0)(\mathbf{1},-\tfrac{1}{3})\,,\\
\omega(\mathbf{8})(\mathbf{1},-\tfrac{1}{3})(\mathbf{1},0) &\leftrightarrow \omega'(\mathbf{8})(\mathbf{1},0)(\mathbf{1},-\tfrac{1}{3})\,,
\end{aligned} \tag{1}$$

while the three generations of ordinary and dark fermions are

$$\begin{aligned}
q_{Li}(\mathbf{3})(\mathbf{2},+\tfrac{1}{6})(\mathbf{1},0) &\leftrightarrow q'_{Ri}(\mathbf{3})(\mathbf{1},0)(\mathbf{2},+\tfrac{1}{6})\,,\\
d_{Ri}(\mathbf{3})(\mathbf{1},-\tfrac{1}{3})(\mathbf{1},0) &\leftrightarrow d'_{Li}(\mathbf{3})(\mathbf{1},0)(\mathbf{1},-\tfrac{1}{3})\,,\\
u_{Ri}(\mathbf{3})(\mathbf{1},+\tfrac{2}{3})(\mathbf{1},0) &\leftrightarrow u'_{Li}(\mathbf{3})(\mathbf{1},0)(\mathbf{1},+\tfrac{2}{3})\,,\\
l_{Li}(\mathbf{1})(\mathbf{2},-\tfrac{1}{2})(\mathbf{1},0) &\leftrightarrow l'_{Ri}(\mathbf{1})(\mathbf{1},0)(\mathbf{2},-\tfrac{1}{2})\,,\\
e_{Ri}(\mathbf{1})(\mathbf{1},-1)(\mathbf{1},0) &\leftrightarrow e'_{Li}(\mathbf{1})(\mathbf{1},0)(\mathbf{1},-1)\,.
\end{aligned} \tag{2}$$

As for the $[SU(2)_L\times SU(2)'_R]$-bidoublet scalar and the three gauge-singlet fermions, they are defined as

$$\Sigma(\mathbf{1})(\mathbf{2},-\frac{1}{2})(\bar{\mathbf{2}},+\frac{1}{2})\leftrightarrow\Sigma^{\dagger}\,,\ \chi_{Ri}(\mathbf{1})(\mathbf{1},0)(\mathbf{1},0)\leftrightarrow\chi^{c}_{Ri}\,. \tag{3}$$

In the above notations, the first, second and third parentheses following the fields are the quantum numbers under the $SU(3)_c$, $SU(2)_L\times U(1)_Y$ and $SU(2)'_R\times U(1)'_Y$ gauge groups, respectively.

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For simplicity, we will not write down the full Lagrangian. Instead, we firstly show the following scalar interactions,

$$\begin{aligned}V \supset\; & \mu_\phi^2\phi^\dagger\phi+\mu_{\phi'}^2\phi'^\dagger\phi'+\mu_\eta^2\eta^*\eta+\mu_{\eta'}^2\eta'^*\eta'+M_\delta^2\delta^\dagger\delta\\&+M_{\delta'}^2\delta'^\dagger\delta'+M_\omega^2\omega^\dagger\omega+M_{\omega'}^2\omega'^\dagger\omega'+M_\Sigma^2\mathrm{Tr}(\Sigma^\dagger\Sigma)\\&+\lambda\delta^\dagger\delta'^\dagger(\omega\eta'+\omega'\eta)+\kappa(\eta^*\omega^\dagger\delta^2+\eta'^*\omega'^\dagger\delta'^2)\\&+\xi\delta^\dagger\delta'\eta\eta'^*+\rho\phi^\dagger\Sigma\phi+\text{H.c.}\,.\end{aligned}\quad(4)$$

where the parity symmetry is assumed to be softly broken, i.e.

$$\mu_\phi^2\neq\mu_{\phi'}^2\,,\quad \mu_\eta^2\neq\mu_{\eta'}^2\,,\quad M_\delta^2\neq M_{\delta'}^2\,,\quad M_\omega^2\neq M_{\omega'}^2\,.\quad(5)$$

Such soft breaking may arise from a spontaneous parity violation [16]. We then give all of the parity-invariant Yukawa couplings,

$$\begin{aligned}\mathcal{L}\supset & -y_\delta(\delta\bar{u}_Re_R^c+\delta'\bar{u}'_Le'^c_L)-y'_\delta(\delta\bar{d}_R\chi_R^c+\delta'\bar{d}'_L\chi_R)\\&-y''_\delta(\delta\bar{q}_Li\tau_2l_L^c+\delta'\bar{q}'_Ri\tau_2l'^c_R)-y_\omega(\omega\bar{u}_R^cd_R+\omega'\bar{u}'^c_Ld'_L)\\&-y'_\omega(\omega\bar{q}_L^ci\tau_2q_L+\omega'\bar{q}'^c_Ri\tau_2q'_R)-y_d(\bar{q}_L\tilde{\phi}d_R+\bar{q}'_R\tilde{\phi}'d'_L)\\&-y_u(\bar{q}_L\phi u_R+\bar{q}'_R\phi'u'_L)-y_e(\bar{l}_L\tilde{\phi}_de_R+\bar{l}'_R\tilde{\phi}'_de'_L)\\&-h(\bar{l}_L\phi\chi_R+\bar{l}'_R\phi'\chi_R^c)-f\bar{l}_L\Sigma l'_R+\text{H.c.}\,,\end{aligned}\quad(6)$$

where a baryon number conservation has been assumed to forbid the other terms involving the leptoquark scalars δ and δ'. We also set the Majorana mass term of the singlet fermions χ_R as below,

$$\mathcal{L}\supset-\frac{1}{2}\mu\bar{\chi}_R^c\chi_R+\text{H.c.}\;\;\text{with}\;\;\mu^T=\mu\,.\quad(7)$$

Furthermore, we introduce a kinetic mixing between the $U(1)_Y$ and $U(1)'_Y$ gauge fields,

$$\mathcal{L}\supset-\frac{\epsilon}{2}B_{\mu\nu}B'^{\mu\nu}\,,\quad(8)$$

where $B_{\mu\nu}$ and $B'^{\mu\nu}$ are the $U(1)_Y$ and $U(1)'_Y$ field strength tensors.

III. SYMMETRY BREAKING

As a result of the softly broken parity, the dark Higgs doublet ϕ' can develop a VEV different from that of the ordinary Higgs doublet ϕ to drive the dark and ordinary electroweak symmetry breaking at different scales,

$$\langle\phi'\rangle\neq\langle\phi\rangle\simeq174\,\text{GeV}\,.\quad(9)$$

Note the VEVs $\langle\phi'\rangle$ and $\langle\phi\rangle$ are both real. After the above symmetry breaking, the heavy Higgs bidoublet Σ can pick up a seesaw-suppressed VEV,

$$\langle\Sigma\rangle\;=\;-\frac{\rho\langle\phi\rangle\langle\phi'\rangle}{M_\Sigma^2}\ll\langle\phi\rangle\,,\langle\phi'\rangle\,.\quad(10)$$

Here we have rotated the cubic coupling ρ to be real. Under the softly broken parity, the dark Higgs singlet η' can acquire a nonzero VEV to spontaneously break the dark electromagnetic symmetry although its ordinary partner η is only allowed to have a zero VEV.

By making a non-unitary transformation [17],

$$\tilde{B}_\mu=B_\mu+\epsilon B'_\mu\,,\quad \tilde{B}'_\mu=\sqrt{1-\epsilon^2}B'_\mu\,,\quad(11)$$

we can remove the $U(1)$ kinetic mixing and then define the following orthogonal fields,

$$\begin{aligned}A_\mu&=W_\mu^3s_W+\tilde{B}_\mu c_W\,,\quad Z_\mu=W_\mu^3c_W-\tilde{B}_\mu s_W\,,\\A'_\mu&=W'^3_\mu s_W+\tilde{B}'_\mu c_W\,,\quad Z'_\mu=W'^3_\mu c_W-\tilde{B}'_\mu s_W\,.\end{aligned}\quad(12)$$

Here $s_W=\sin\theta_W$ and $c_W=\cos\theta_W$ with θ_W being the Weinberg angle, while W^3 and W'^3 are the diagonal components of the $SU(2)_L$ and $SU(2)'_L$ gauge fields. In the above orthogonal base, the field A is exactly massless and is the ordinary photon, according to the unbroken electromagnetic symmetry in the ordinary sector, while the dark photon A' mixes with the Z and Z' bosons. As the dark electromagnetic symmetry now is broken by the VEV $\langle\eta'\rangle$, the dark photon can have a mass

$$\begin{aligned}m_{A'}\;&\simeq\;\frac{2\sqrt{2\pi\alpha}}{3}\langle\eta'\rangle\simeq1\,\text{GeV}\left(\frac{\langle\eta'\rangle}{7\,\text{GeV}}\right)\\&\quad\text{for}\;\;\epsilon\ll1\,,\quad\langle\eta'\rangle\ll\langle\phi'\rangle\,,\end{aligned}\quad(13)$$

with $\alpha=e^2/(4\pi)\simeq1/137$ being the fine structure constant. The dark photon can couple to the ordinary fermions besides the dark fermions [18],

$$\begin{aligned}\mathcal{L}\;\supset\;& eA'_\mu\left\{\frac{\epsilon}{4}\left[\bar{e}\gamma^\mu\left(3+\gamma_5\right)e+\bar{\nu}\gamma^\mu(1-\gamma_5)\nu\right.\right.\\&\left.+\bar{d}\gamma^\mu\left(\frac{1}{3}+\gamma_5\right)d-\bar{u}\gamma^\mu\left(\frac{5}{3}+\gamma_5\right)u\right]\\&\left.+\left(-\frac{1}{3}\bar{d}'\gamma^\mu d'+\frac{2}{3}\bar{u}'\gamma^\mu u'-\bar{e}'\gamma^\mu e'\right)\right\}\,.\end{aligned}\quad(14)$$

As long as the dark photon A' is heavy enough, it can efficiently decay into the ordinary fermions. The Higgs boson from the dark Higgs scalar η' can decay into the dark photon. The dark scalar η can have a three-body decay mode into the ordinary leptoquark and diquark scalars (two δ and one ω) or a two-body decay mode into the ordinary and dark leptoquark scalars (one δ and one δ').

IV. PREDICTIVE DARK MATTER MASS AND ITS IMPLICATION

The Yukawa couplings (6) can tell us the masses of the dark quarks and charged leptons from the ordinary ones,

$$\begin{aligned}&\frac{m_{d'}}{m_d}=\frac{m_{s'}}{m_s}=\frac{m_{b'}}{m_b}=\frac{m_{u'}}{m_u}=\frac{m_{c'}}{m_c}=\frac{m_{t'}}{m_t}=\frac{m_{e'}}{m_e}\\&=\frac{m_{\mu'}}{m_\mu}=\frac{m_{\tau'}}{m_\tau}=\frac{\langle\phi'\rangle}{\langle\phi\rangle}\,.\end{aligned}\quad(15)$$

As we will show later, the dark neutrinos and quarks are unstable and don't contribute to the relic density of the present universe. In other words, the dark electron is the unique stable particle in the dark sector. The dark electron can annihilate into the dark photon,

$$\sigma_{e'} = \langle \sigma_{e'^+ e'^- \to A' A'} v_{\rm rel} \rangle = \frac{\pi \alpha^2}{m_{e'}^2} \quad \text{for} \quad m_{e'} \gg m_{A'} \,. \quad (16)$$

The relic density of the dark electron then can be calculated by [21]

$$\Omega_{e'} h^2 = \frac{1.07 \times 10^9 m_{e'}}{\sqrt{g_*} M_{\rm Pl} \sigma_{e'} T_f({\rm GeV})} \,, \quad (17)$$

where $M_{\rm Pl} \simeq 1.22 \times 10^{19}\,\text{GeV}$ is the Planck mass, T_f is the freeze-out temperature [21],

$$\begin{aligned} \frac{m_{e'}}{T_f} &= \ln(0.152\, M_{\rm Pl} m_{e'} \sigma_{e'}/\sqrt{g_*}) \\ &\quad -\frac{1}{2} \ln[\ln(0.152\, M_{\rm Pl} m_{e'} \sigma_{e'}/\sqrt{g_*})] \,, \end{aligned} \quad (18)$$

while $g_* = g_*(T)$ is the number of relativistic degrees of freedom. Remarkably, the relic density (17) only depends on one unknown parameter $m_{e'}$, the dark matter mass, similar to that in some minimal dark matter scenarios where the exotic dark matter particles only have the SM gauge interactions [22]. By inputting $g_* = 90$ [21], we find if the dark electron is expected to account for the dark matter relic density [23], its mass should have the following value,

$$m_{e'} \simeq 302.1 \pm 3.5\,\text{GeV for } \Omega_{e'} h^2 = 0.1199 \pm 0.0027 \quad (19)$$

Through the exchange of the dark photon, the dark electron can scatter off the ordinary nucleon,

$$\begin{aligned} \sigma_{e'N \to e'N} &\simeq \epsilon^2 \frac{\pi \alpha^2 \mu_r^2}{m_{A'}^4} \left[\frac{3Z + (A-Z)}{A}\right]^2 , \\ &\simeq 10^{-45}\,\text{cm}^2 \left(\frac{\epsilon}{1.25 \times 10^{-7}}\right)^2 \left(\frac{\mu_r}{1\,\text{GeV}}\right)^2 \\ &\quad \times \left(\frac{1\,\text{GeV}}{m_{A'}}\right)^4 \left[\frac{3Z + (A-Z)}{A}\right]^2 , \end{aligned} \quad (20)$$

which is accessible to the ongoing and future dark matter direct detection experiments such as the XENON100 [24] and XENON1T experiments. Here Z and $A - Z$ are the numbers of proton and neutron within the target nucleus, while

$$\mu_r = \frac{m_{e'} m_N}{m_{e'} + m_N} \simeq m_N \simeq 1\,\text{GeV} \quad \text{for} \quad m_{e'} \simeq 302\,\text{GeV} \,, \quad (21)$$

is the reduced mass. The dark photon will also mediate a self-interaction of the dark electron. For example, we can have a self-interacting cross section,

$$\begin{aligned} \frac{\sigma_{e'^+ e'^- \to e'^+ e'^-}}{m_{e'}} &\simeq \frac{4\pi \alpha^2 m_{e'}}{m_{A'}^4} \\ &\simeq 1.6 \times 10^{-42}\,\text{cm}^3 \left(\frac{m_{e'}}{302\,\text{GeV}}\right) \left(\frac{1\,\text{GeV}}{m_{A'}}\right)^4 , \end{aligned} \quad (22)$$

which is easy to satisfy the observation limits [25].

As the dark electron mass is determined, we can fix the dark VEV $\langle \phi' \rangle$ from Eq. (15),

$$\begin{aligned} \langle \phi' \rangle &= \frac{m_{e'}}{m_e} \langle \phi \rangle \simeq 10^8\,\text{GeV} \\ &\quad \times \left(\frac{m_{e'}}{302\,\text{GeV}}\right) \left(\frac{0.511\,\text{MeV}}{m_e}\right) \left(\frac{\langle \phi \rangle}{174\,\text{GeV}}\right) . \end{aligned} \quad (23)$$

Accordingly, the masses of the other dark quarks and charged leptons should be

$$\begin{array}{lclcl} m_{d'} &=& 2.8\,\text{TeV} & \text{for } m_d &=& 4.8\,\text{MeV}\,, \\ m_{u'} &=& 1.3\,\text{TeV} & \text{for } m_u &=& 2.3\,\text{MeV}\,, \\ m_{s'} &=& 55\,\text{TeV} & \text{for } m_s &=& 95\,\text{MeV}\,, \\ m_{c'} &=& 738.6\,\text{TeV} & \text{for } m_c &=& 1.275\,\text{GeV}\,, \\ m_{b'} &=& 2.42 \times 10^3\,\text{TeV} & \text{for } m_b &=& 4.18\,\text{GeV}\,, \\ m_{t'} &=& 9.971 \times 10^5\,\text{TeV} & \text{for } m_t &=& 173.5\,\text{GeV}\,, \\ m_{\mu'} &=& 61.23\,\text{TeV} & \text{for } m_\mu &=& 105.7\,\text{MeV}\,, \\ m_{\tau'} &=& 1.029 \times 10^3\,\text{TeV} & \text{for } m_\tau &=& 1.777\,\text{GeV}\,. \end{array} \quad (24)$$

V. CONSEQUENCE AND FATE OF DARK QUARKS

Not only the ordinary quarks but also the dark quarks are related to the non-perturbative QCD Lagrangian,

$$\mathcal{L} \supset \bar{\theta} \frac{g_3^2}{32\pi^2} G\tilde{G} \quad \text{with} \quad \bar{\theta} = \theta + \text{ArgDet}(M_u M_d) \,, \quad (25)$$

where θ is the original QCD phase, while M_u and M_d are the mass matrices of the up-type and down-type quarks,

$$\begin{aligned} \mathcal{L} &\supset -[\bar{d}_L \; \bar{d}'_L] M_d \begin{bmatrix} d_R \\ d'_R \end{bmatrix} - [\bar{u}_L \; \bar{u}'_L] M_u \begin{bmatrix} u_R \\ u'_R \end{bmatrix} + \text{H.c.} \\ &\quad \text{with} \;\; M_{d(u)} = \begin{bmatrix} y_{d(u)} \langle \phi \rangle & 0 \\ 0 & y^\dagger_{d(u)} \langle \phi' \rangle \end{bmatrix} . \end{aligned} \quad (26)$$

Obviously, the θ-term in the QCD Lagrangian should be zero because of the parity invariance while the [ArgDet$(M_u M_d)$]-term should be also trivial due to the real determinants Det(M_d) and Det(M_u). We hence can obtain a vanishing strong CP phase $\bar{\theta} = 0$ [7, 19, 20].

From Eqs. (4) and (6), it is easy to see that after the dark electromagnetic symmetry breaking, a dark up quark can decay into a dark electron and four ordinary fermions through the mediations of the real and/or virtual colored scalars. For example, we can naively esti-

mate

$$\begin{aligned}
\Gamma_{u'\to e'^+ e^+ \bar{u}\bar{u}\bar{d}} &\sim \frac{1}{2^{16}\pi^7}\frac{\langle\eta'\rangle^2 m_{u'}^{11}}{M_{\delta'}^4 M_\delta^4 M_\omega^4}|\lambda|^2[(y_\delta^\dagger y_\delta)_{11}\\
&\quad +(y_\delta''^\dagger y_\delta'')_{11}][\mathrm{Tr}(y_\delta^\dagger y_\delta)+\mathrm{Tr}(y_\delta''^\dagger y_\delta'')]\\
&\quad \times[\mathrm{Tr}(y_\omega^\dagger y_\omega)+2\mathrm{Tr}(y_\omega'^\dagger y_\omega')]\\
&\simeq \frac{1}{3.5\times 10^{-11}\,\mathrm{sec}}\left(\frac{\langle\eta\rangle}{7\,\mathrm{GeV}}\right)^2\left(\frac{m_{u'}}{1.3\,\mathrm{TeV}}\right)^{11}\\
&\quad \times\left(\frac{10\,m_{u'}}{M_{\delta'}}\right)^4\left(\frac{m_{u'}}{M_\delta}\right)^4\left(\frac{m_{u'}}{M_\omega}\right)^4\\
&\quad \times|\lambda|^2[(y_\delta^\dagger y_\delta)_{11}+(y_\delta''^\dagger y_\delta'')_{11}][\mathrm{Tr}(y_\delta^\dagger y_\delta)\\
&\quad +\mathrm{Tr}(y_\delta''^\dagger y_\delta'')][\mathrm{Tr}(y_\omega^\dagger y_\omega)+2\mathrm{Tr}(y_\omega'^\dagger y_\omega')]\,.
\end{aligned} \tag{27}$$

So, the lightest dark quark can have a very short lifetime for a proper parameter choice.

VI. NEUTRINO MASSES AND BARYON ASYMMETRY

From the Yukawa couplings (6), we can read the mass terms involving the ordinary neutrinos ν_L, the dark neutrinos ν'_R and the singlet fermions χ_R,

$$\mathcal{L} \supset -\frac{1}{2}[\,\bar{\nu}_L\ \ \bar{\nu}_R'^c\ \ \bar{\chi}_R^c\,]\begin{bmatrix} 0 & f\langle\Sigma\rangle & h\langle\phi_u\rangle \\ f^T\langle\Sigma\rangle & 0 & h^*\langle\phi'\rangle \\ h^T\langle\phi\rangle & h^\dagger\langle\phi'\rangle & \mu\end{bmatrix}\begin{bmatrix}\nu_L^c\\ \nu'_R\\ \chi_R\end{bmatrix} +\mathrm{H.c.}\,. \tag{28}$$

Here the Yukawa coupling matrix f are hermitian due to the parity symmetry. For $h\langle\phi'\rangle \gg \mu, h\langle\phi\rangle, f\langle\Sigma\rangle$, the seesaw mechanism can be applied to give the masses of the ordinary neutrinos,

$$\begin{aligned}
m_\nu &= f\langle\Sigma\rangle\frac{1}{h^\dagger\langle\phi'\rangle}\mu\frac{1}{h^*\langle\phi'\rangle}f^T\langle\Sigma\rangle\\
&\quad -(f\langle\Sigma\rangle\frac{1}{h^\dagger}h^T+h\frac{1}{h^*}f^T\langle\Sigma\rangle)\frac{\langle\phi\rangle}{\langle\phi'\rangle}\\
&= \tilde{f}\frac{1}{\hat{h}}\tilde{\mu}\frac{1}{\hat{h}}\tilde{f}^T\frac{\langle\Sigma\rangle^2}{\langle\phi'\rangle^2}-(\tilde{f}+\tilde{f}^T)\frac{\langle\phi\rangle\langle\Sigma\rangle}{\langle\phi'\rangle}\,,
\end{aligned} \tag{29}$$

where we have defined

$$\begin{aligned}
&\hat{h} = U_R h V_R^T = \mathrm{diag}\{\hat{h}_1\,,\ \hat{h}_2\,,\ \hat{h}_3\}\,,\\
&\tilde{f} = f U_R^\dagger\,,\ \tilde{\mu} = V_R\mu V_R^T\,.
\end{aligned} \tag{30}$$

The first term of Eq. (29) is the inverse seesaw [12], while the second term is the linear seesaw [13].

In this inverse and linear seesaw scenario, a dark neutrino ν'_{Ri} and a singlet fermion χ_{Ri} actually form two quasi-degenerate Majorana fermions [26],

$$\begin{aligned}
N_i^+ &\simeq \frac{1}{\sqrt{2}}(\nu'_{Ri}+\chi_{Ri}+\nu_{Ri}'^c+\chi_{Ri}^c)\ \ \text{with}\\
&\quad M_{N_i^+} = \hat{h}_i\langle\phi'\rangle+\frac{1}{2}\tilde{\mu}_{ii}\,,\\
N_i^- &\simeq \frac{i}{\sqrt{2}}(\nu'_{Ri}-\chi_{Ri}-\nu_{Ri}'^c+\chi_{Ri}^c)\ \ \text{with}\\
&\quad M_{N_i^-} = \hat{h}_i\langle\phi'\rangle-\frac{1}{2}\tilde{\mu}_{ii}\,.
\end{aligned} \tag{31}$$

The Yukawa couplings of the quasi-degenerate Majorana fermions to the ordinary leptons and Higgs scalar should be

$$\begin{aligned}
\mathcal{L} &\supset -y_+\bar{l}_L\phi N^+ - y_-\bar{l}_L\phi N^- +\mathrm{H.c.}\ \ \text{with}\\
&\quad y_+ = \frac{1}{\sqrt{2}}(U_R^\dagger\hat{h}-\frac{\rho\langle\phi'\rangle}{M_\Sigma^2}\tilde{f}) = \frac{1}{\sqrt{2}}(U_R^\dagger\hat{h}+\bar{f})\,,\\
&\quad y_- = \frac{i}{\sqrt{2}}(U_R^\dagger\hat{h}+\frac{\rho\langle\phi'\rangle}{M_\Sigma^2}\tilde{f}) = \frac{i}{\sqrt{2}}(U_R^\dagger\hat{h}-\bar{f})\,.
\end{aligned} \tag{32}$$

Such quasi-degenerate Majorana fermions can accommodate a resonant leptogenesis to generate the baryon asymmetry in the universe. Following [9], we compute the CP asymmetries from self-energy corrections,

$$\begin{aligned}
\varepsilon_{N_i^\pm} &= \frac{\Gamma(N_i^\pm\to l_L\bar{\phi})-\Gamma(N_i^\pm\to\bar{l}_L\phi)}{\Gamma(N_i^\pm\to l_L\bar{\phi})+\Gamma(N_i^\pm\to\bar{l}_L\phi)}\\
&\simeq \frac{\mathrm{Im}\{[(y_+^\dagger y_-)_{ii}]^2\}}{8\pi(y_\pm^\dagger y_\pm)_{ii}}\frac{r_{N_i}}{r_{N_i}^2+\frac{1}{64\pi^2}[(y_\mp^\dagger y_\mp)_{ii}]^2}\,,
\end{aligned} \tag{33}$$

which can be specified by expanding

$$\begin{aligned}
r_{N_i} &= \frac{M_{N_i^+}^2-M_{N_i^-}^2}{M_{N_i^+}M_{N_i^-}} \simeq \frac{2\tilde{\mu}_{ii}}{\hat{h}_i\langle\phi'\rangle}\,,\\
(y_\pm^\dagger y_\pm)_{ii} &= \frac{1}{2}\{\hat{h}_i^2\pm 2\hat{h}_i\mathrm{Re}[(U_R\bar{f})_{ii}]+(\bar{f}^\dagger\bar{f})_{ii}\}\,,\\
\mathrm{Im}\{[(y_\pm^\dagger y_\pm)_{ii}]^2\} &= \hat{h}_i[\hat{h}_i^2-(\bar{f}^\dagger\bar{f})_{ii}]\mathrm{Im}[(U_R\bar{f})_{ii}]\,.
\end{aligned} \tag{34}$$

When the Majorana fermions $N_i^\pm$ go out of equilibrium, their CP-violating decays can generate a lepton asymmetry in the ordinary leptons l_L. This lepton asymmetry then will be partially converted to a baryon asymmetry through the sphaleron processes [27]. The induced baryon asymmetry can be approximately described by [21],

$$\begin{aligned}
\eta_B &\simeq -\frac{28}{79}\times\sum\frac{\varepsilon_{N_i^\pm}\kappa_{N_i^\pm}}{g_*}\ \ \text{with}\\
&\quad \kappa_{N_i^\pm} \simeq \begin{cases} 1 & \text{for}\ \ K_{N_i^\pm}\ll 1\,,\\ \frac{0.3}{K_{N_i^\pm}(\ln K_{N_i})^{0.6}} & \text{for}\ \ K_{N_i^\pm}\gtrsim 1\,,\end{cases}
\end{aligned} \tag{35}$$

where the parameters $K_{N_i^\pm}$ are defined as

$$K_{N_i^\pm} = \frac{\Gamma_{N_i^\pm}}{2H(T)}\Big|_{T=M_{N_i^\pm}}\,, \tag{36}$$

with $\Gamma_{N_i^\pm}$ being the decay width,

$$\Gamma_{N_i^\pm} = \Gamma(N_i^\pm \to l_L\bar{\phi}) + \Gamma(N_i^\pm \to \bar{l}_L\phi) = \frac{1}{8\pi}(y_\pm^\dagger y_\pm)_{ii} M_{N_i^\pm}\,, \quad (37)$$

and $H(T)$ being the Hubble constant,

$$H(T) = \left(\frac{8\pi^3 g_*}{90}\right)^{\frac{1}{2}} \frac{T^2}{M_{\rm Pl}}\,. \quad (38)$$

We now show a proper parameter choice can simultaneously result in the required neutrino masses and baryon asymmetry. For this purpose, we set

$$M_\Sigma = \langle\phi'\rangle = 10^8\,\mathrm{GeV}\,, \quad \rho = \langle\phi\rangle\,, \quad (39)$$

to give the seesaw-suppressed VEV in Eq. (10) and the effective Yukawa couplings in Eq. (32),

$$\langle\Sigma\rangle \simeq -0.3\,\mathrm{MeV}\,, \quad \bar{f} \simeq -1.74\times 10^{-6}\,\tilde{f}\,. \quad (40)$$

By further inputting,

$$\hat{h}_1 = 10^{-5} \ll \hat{h}_{2,3}\,, \quad \tilde{\mu}_{ij} = 10\,\mathrm{eV}\,, \quad \tilde{f} = \mathcal{O}(0.1-1)\,, \quad (41)$$

the neutrino masses (29) can be dominated by the linear seesaw,

$$m_\nu \simeq -(\tilde{f}+\tilde{f}^T)\frac{\langle\phi\rangle\langle\Sigma\rangle}{\langle\phi'\rangle} = 0.5\,\mathrm{eV}(\tilde{f}+\tilde{f}^T)\,. \quad (42)$$

The above parameter choice can also induce

$$M_{N_1^\pm} = 1\,\mathrm{TeV} \ll M_{N_{2,3}^\pm}\,, \quad K_{N_1^\pm} \simeq 700\left(\frac{110.75}{g_*}\right)^{\frac{1}{2}}\,,$$
$$\varepsilon_{N_1^\pm} \simeq 0.98\times 10^{-3}\left(\frac{\mathrm{Im}[(U_R\tilde{f})_{11}]}{-0.0154}\right)\,, \quad (43)$$

to explain the measured baryon asymmetry [23],

$$\eta_B \simeq 3.81\times 10^{-9}\times(0.02205\pm 0.00028) \simeq (0.8401\pm 0.0107)\times 10^{-10}\,. \quad (44)$$

Note the dark leptoquark scalar and gauge bosons can mediate some scattering and annihilating processes of the decaying Majorana fermions. As an example, we check the processes $N_1^\pm N_1^\pm \to d'\bar{d}'$ mediated by the dark leptoquark δ' and find,

$$\Gamma_S \sim [(y_\delta'^\dagger y_\delta')_{11}]^2\frac{T^5}{M_{\delta'}^4} \sim H(T) \Rightarrow$$
$$T \sim 3.7\,\mathrm{TeV}\left[\frac{10^{-6}}{(y_\delta'^\dagger y_\delta')_{11}}\right]^{\frac{2}{3}}\left(\frac{M_{\delta'}}{10\,m_{e'}}\right)^{\frac{4}{3}}\left(\frac{g_*}{159}\right)^{\frac{1}{6}} \quad (45)$$

Similarly, the other scattering and annihilating processes can also decouple at a temperature above the leptogenesis epoch [28].

VII. CONCLUSION

In the presence of an $SU(2)_R'\times U(1)_Y'$ dark parity extension of the $SU(2)_L\times U(1)_Y$ ordinary electroweak theory, the strong CP problem can be solved without resorting to an axion. In this framework, we consider three gauge-singlet fermions with small Majorana masses and an $[SU(2)_L\times SU(2)_R']$-bidoublet Higgs scalar with seesaw-suppressed VEV to generate the baryon asymmetry by resonant leptogenesis and the neutrino masses by inverse and linear seesaw. We also introduce some colored scalars to mediate a fast decay of the lightest dark quark after the dark electromagnetic symmetry breaking. The dark electron can keep stable to account for the dark matter relic density if it has a determined mass around 302 GeV. Benefited from the $U(1)_Y\times U(1)_Y'$ kinetic mixing, the dark matter particle can be verified by the dark matter direct detection experiments .

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